Supersymmetric Dark Matter: Relic Density and Detection[*]

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Abstract

The Lightest Supersymmetric Particle (LSP) makes a good Dark Matter (DM) candidate, since its relic density quite naturally comes out close to the cosmologically required value. This is true even in minimal Supergravity models with radiative symmetry breaking, which have a rather small number of free parameters. On the other hand, the experimental detection of SUSY DM might be quite difficult, necessitating km^2 size neutrino detectors or direct detection experiments with several tons of detector material.

1) Introduction

The possibility that all known particles have superpartners [1] with mass ~ 1 TeV or less is now being taken quite seriously. The main motivation for the introduction of supersymmetry is that it makes the theory technically natural, i.e. protects the weak scale against large radiative corrections [2]. It has recently been found [3] that supersymmetry also facilitates the introduction of a Grand Unified gauge group at scale $M_X \simeq 10^{16}$ GeV. Of more immediate interest for this talk is that in the simplest, "R–parity invariant" SUSY models the lightest superparticle (LSP) is absolutely stable. This means that some of the LSPs produced just after the Big Bang should still be around today. Searches for exotic isotopes [4] put very stringent upper limits on the present abundance of charged and/or strongly interacting particles; if its density is to be cosmologically significant the LSP therefore has to be neutral.

Fortunately in many models the LSP more or less automatically comes out to be the lightest neutralino, which satisfies this requirement. Assuming minimal particle content the LSP is then a mixture of four different current eigenstates: The bino $\tilde{B}$, the neutral wino $\tilde{W}$, and the two higgsinos $\tilde{h}_1^0$, $\tilde{h}_2^0$. The mixing is governed by the following mass matrix:

$$\mathcal{M}^0 = \begin{pmatrix} M_1 & 0 & -M_Z\sin\theta_W\cos\beta & M_Z\sin\theta_W\sin\beta \\ 0 & M_2 & M_Z\cos\theta_W\cos\beta & -M_Z\cos\theta_W\sin\beta \\ -M_Z\sin\theta_W\cos\beta & M_Z\cos\theta_W\cos\beta & 0 & -\mu \\ M_Z\sin\theta_W\sin\beta & -M_Z\cos\theta_W\sin\beta & -\mu & 0 \end{pmatrix}.$$

(1)

Here, M_1 and M_2 are SUSY breaking $U(1)$ and $SU(2)$ gaugino masses, μ is the supersymmetric Higgs(ino) mass and $\tan\beta \equiv \langle H_2^0 \rangle / \langle H_1^0 \rangle$ is the ratio of vevs. For the remainder of this talk I will assume gaugino masses to be unified at scale M_X, leading to the following relation at the weak scale:

$$M_1 = \frac{5}{3}\tan^2\theta_W M_2 \simeq 0.5M_2,$$

(2)

at 1–loop order [1]. This then implies that there are basically only three possibilities for the LSP:

- If $|\mu| \geq |M_2|$ the LSP is gaugino–like (mostly photino for small $|M_1|$, mostly bino for $|M_1| \geq M_Z$); its mass is $\sim M_1$.

- If $|\mu| \simeq |M_1|$ the LSP is a mixed state, i.e. both higgsino and gaugino components will be sizable.

- If $|\mu| < |M_1|$ the LSP is higgsino–like with mass $\sim \mu$. Notice that in this case the next–to–lightest sparticle, either the second neutralino or the lighter chargino, will be close in mass to the LSP.

We will see in the next section that only one of these three types of neutralinos makes a good particle DM candidate.

2) LSP relic abundance

By now there are quite a few calculations or LSP relic abundances, stretching back more than 10 years [5]. There is general agreement that, over wide regions of parameter space, the LSP does make a good particle DM candidate. Here I will briefly summarize the results of ref.[6], which included all 2–body final states accessible to LSP annihilation at tree level, and which was among the first papers to use a spartcicle spectrum as predicted in minimal supergravity models with radiative gauge symmetry breaking [7].

In the very early Universe LSPs were in thermal equilibrium with ordinary particles. However, as the Universe expanded and cooled the rate of reactions involving LSPs eventually dropped below the expansion rate. At this point the LSP density "froze out", i.e. it basically remained constant (over a co–moving volume). The larger the LSP annihilation cross section, the smaller the freeze–out temperature and hence the smaller the LSP relic abundance (due to the Boltzmann factor). Freeze–out approximately occurs at $T_F \simeq m_\chi/20$, i.e. LSPs are nonrelativistic at freeze–out (hence "cold DM"). It is therefore usually (but not always [8]) sufficient to expand the LSP annihilation cross section in powers of the relative cms velocity v:

$$v\sigma_{\mathrm{ann}}(\chi\chi \rightarrow \mathrm{anything}) = a + bv^2 + \mathcal{O}(v^4). \tag{3}$$

The present relic density is then very roughly given by

$$\Omega_\chi h^2 \simeq \frac{0.08\mathrm{pb}}{a + b/7}. \tag{4}$$

Here Ω_χ is the LSP relic density in units of the critical (closure) density $\rho_c \simeq 2 \cdot 10^{-29}$ g/cm^3, and h is the Hubble expansion parameter in units of 100 km/(sec·Mpc).

Note that all final states that are kinematically accessible contribute in eq.(3). For each final state there are usually contributions from $t-$channel diagrams (where a sfermion, neutralino or chargino is exchanged) as well as $s-$channel diagrams (where a Z or Higgs boson is exchanged). Ω_χ thus depends on the *entire* (s)particle spectrum. A more or less exhaustive scan of parameter space is therefore only possible in models where supersymmetry breaking is described by a small number of free parameters, as compared to the dozens of parameters in general softly broken supersymmetry. One particularly attractive class are minimal supergravity (mSUGRA) models [7]. Here

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one assumes one common nonsupersymmetric squared scalar mass m^2 (sometimes called m_0^2) and one common gaugino mass M (sometimes called $m_{1/2}$), as well as one common trilinear soft breaking parameter A. The sparticle spectrum is assumed to have this very simple form only at ultrahigh energies near the Planck scale. Radiative corrections, most conveniently described by a set of renormalization group equations [9], change the spectrum at lower energies. In particular, corrections involving the Yukawa coupling of the top quark can drive the squared mass of one of the Higgs bosons to negative values, thereby triggering electroweak gauge symmetry breaking.

In ref.[6] we studied the LSP relic density in such mSUGRA models*. Fig.1, taken from that paper, shows contours of $\Omega_\chi h^2$ in the $(M, \tan\beta)$ plane. We see that a good part of the experimentally and theoretically allowed region (within the outer, dotted lines) satisfies $0.25 \leq \Omega_\chi h^2 \leq 1$, i.e. allows for a flat Universe as favored by inflationary scenarios for the range $0.5 \leq h \leq 1$ currently favored by cosmologists (area between the long dashed and solid lines). The short dashed lines show locations of $s-$channel poles, i.e. $2m_\chi = m_{Z,h,P}$, where the relic density is usually very small.

From this and similar figures one concludes that $\Omega_\chi h^2$ indeed "naturally" comes out in the right ball park if m and M are of the order of a few hundred GeV, just in the range expected from naturalness arguments. Unfortunately the requirement $\Omega_\chi h^2 \leq 1$ does not translate into a strict upper bound on sparticle masses, however. The reason is that LSP annihilation can be greatly enhanced if $2m_\chi$ is close to the mass of the pseudoscalar boson P (unlike Z and h exchange, P exchange gives an $S-$wave pole, i.e. contributes at order v^0). In mSUGRA this requires large $\tan\beta$, since one needs a large b Yukawa coupling to reduce m_P. This loophole allows cosmologically viable SUSY breaking at scales well above those favored by naturalness arguments.

Another important result is that a cosmologically interesting LSP is almost always gaugino–like. This statement is true even in more general SUSY models[5]. Higgsino–like and mixed LSPs have large couplings to gauge and Higgs bosons; their density therefore drops rapidly for $m_\chi > m_W$, reaching interesting values again only for $m_\chi > 0.5$ TeV, which is already uncomfortably heavy for the lightest superparticle if SUSY is to stabilize the gauge hierarchy. If $m_\chi < m_W$ their density is suppressed by the proximity of $s-$channel poles (to make this argument watertight one has to include co–annihilation of the LSP with the next–to–lightest sparticle [10]). In mSUGRA $|\mu|$ has to be increased along with m_t to give the proper W and Z boson masses, so that for $m_t \geq 150$ GeV the LSP is almost always gaugino–like; this can be regarded as one of the successes of such models.

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*In this paper a relation between the SUSY breaking parameters A and B was assumed, but relaxing this assumption does not alter the overall conclusions.

3) Search for LSP annihilation in the Sun and Earth

How would one go about looking experimentally for the supersymmetric DM particles that are predicted to surround us? One idea [11] is to look for energetic muon neutrinos emerging from the centre of the Earth or Sun. DM particles have a finite chance to interact with nuclei in celestial bodies, thereby losing enough energy to become trapped gravitationally. In subesquent interactions those particles lose even more energy, and finally they become concentrated in the centre of these bodies. After some time so many LSP particles should have accumulated in the centre of the Sun and Earth that they begin to annihilate at a significant rate; capture and annihilation eventually reach equilibrium[†]. Some fraction of annihilation events will give rise to energetic muon neutrinos, which can be detected in nucleon decay detectors like Kamiokande [12] or in dedicated neutrino detectors like AMANDA [13]. The total signal rate is proportional to

$$\text{Signal} \propto (\text{capture rate}) \cdot \sum_X Br(\chi\chi \to X) \int_{E_{thr}}^{m_\chi} dE_\nu (E_\nu)^2 \frac{d\Gamma}{dE_\nu}. \tag{5}$$

Since we assume dynamical equilibrium between capture and annihilation the overall size of the signal can be expressed in terms of the capture rate. However, not all final states X contribute equally to the signal; rather, it is proportional to the third moment of the neutrino energy spectrum characteristic for that final state. (One factor of E_ν appears since the $\nu \to \mu$ cross section grows with energy, the second is due to the increased range of the produced μ.) Present searches already exclude [12] some combinations of parameters; however, calculations indicate [14] that one needs a detector with effective area of at least 1 km^2 to cover most of the interesting parameter space.

   Previous calculations [11, 14] only included final states X accessible at tree level. Especially for gaugino–like LSPs with mass below m_t the process

$$\chi\chi \to gg \tag{6}$$

can also be relevant [15]. The reason is that LSP annihilation even in the Sun occurs almost at rest; hence only the $\mathcal{O}(v^0)$ term a in the annihilation cross section (3) is relevant. For $f\bar{f}$ final states this term is proportional to m_f^2. LSPs at rest therefore predominantly annihilate into the most massive accessible fermions (c or b quarks or τ leptons). This is fortunate, since their (semi–)leptonic decays can give rise to hard neutrinos. However, massless quarks do contribute to the process (6) at the 1–loop level. The relative size of the corresponding cross section is therefore roughly (for bino–like LSP)

$$\frac{\sigma(\chi\chi \to gg)}{\sigma(\chi\chi \to b\bar{b})} \sim \left(\frac{\alpha_s}{\pi}\right)^2 \frac{\left(\sum_q Y_q^2\right)^2}{Y_b^4} \left(\frac{m_\chi}{m_b}\right)^2, \tag{7}$$

[†]In case of the Sun equilibrium should have been reached long ago; for the Earth this is only true if LSPs are not too heavy [12]

which can exceed unity. Unfortunately gluons do not produce hard neutrinos; the process (6) therefore reduces the signal by lowering the branching ratio for detectable processes. Fig.2 shows that the reduction can amount to a factor of 2 if LSPs are indeed gaugino–like and not much lighter than the top quark. QCD corrections can therefore make the search for neutrino signals from LSP annihilation significantly more difficult, but they should not make it impossible.

4) Direct LSP detection

Another possibility to test the existence of SUSY DM is to search for the elastic scattering of relic LSPs off the nuclei in a detector [16]. In order to estimate the expected event rate one obviously needs to know the LSP–nucleus scattering cross section. As a first step one has to compute the S matrix element for scattering off a single nucleon. In general one distinguishes two kinds of LSP–nucleon interactions: Those that depend on the spin of the nucleon and those that don't. The former interactions arise from Z boson and $\tilde{q}$ exchange while the latter are due to scalar Higgs boson and $\tilde{q}$ exchange [17]. This last contribution – the spin–independent $\tilde{q}$–exchange term – involves some subtleties, as shown in refs.[18]. To see that we start from the LSP–quark–squark interaction

$$\mathcal{L}_{\chi q\tilde{q}} = \bar{q}\sum_{i=1}^{2}\left(a_{\tilde{q}_i} + b_{\tilde{q}_i}\gamma_5\right)\chi\tilde{q}_i + h.c., \tag{8}$$

where the couplings $a_{\tilde{q}_i}$, $b_{\tilde{q}_i}$ depend on the mixing angles both in the neutralino sector and between the two squark eigenstates of a given flavor $\tilde{q}_i$. Eq.(8) allows to compute an effective LSP–quark interaction by integrating out the squark fields; this is a good approximation since one is interested in momentum transfers (much) less than 1 GeV, well below the squark mass. In leading order in inverse squark mass one then gets a spin–independent interaction $\propto a_{\tilde{q}_i}^2 - b_{\tilde{q}_i}^2$, which is nonzero only if chirality is broken in the (s)quark sector. In the Standard Model and its supersymmetric extension chirality is broken only by terms involving the quark mass, so that the leading (in powers of $m_{\tilde{q}}^{-2}$) spin–independent LSP–quark interaction is $\propto m_q$. (This is obviously also true for the Higgs exchange contribution.)

This special role played by massive quarks was first recognized by Griest [17]. He and subsequent authors estimated the heavy quark contribution to LSP–nucleon scattering using a trick due to Shifman et al.[19]: By closing the heavy quark line in a loop and attaching two gluons the matrix element $\langle N|m_q\bar{q}q|N\rangle$ can (for $m_q \gg \Lambda_{\mathrm{QCD}}$) be related to the matrix element $\langle N|F_{\mu\nu}F^{\mu\nu}|N\rangle$ which in turn is related to the nucleon mass. Shifman et al. had used this trick in an estimate of Higgs–nucleon scattering rates. However, in the present case the LSP–quark interaction itself is due to the exchange of a strongly interacting particle, a squark; by using the SVZ trick one effectively contracts the $\tilde{q}$ propagator to a point *inside* a loop, see fig.3. This procedure

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cannot be expected to produce reliable answers when m_q and $m_{\tilde{q}}$ are comparable (which might be true for the top quark); it will also fail to reproduce higher order (in $m_{\tilde{q}}^{-2}$) terms.

We therefore computed [18] the full LSP–gluon scattering amplitude. Up to permutations of external lines there are four different Feynman diagrams, only one of which is contained in the effective Lagrangian treatment oulined above. This calculation reproduced the $(a_{\tilde{q}_i}^2 - b_{\tilde{q}_i}^2) F_{\mu\nu} F^{\mu\nu} \bar{\chi}\chi$ term found previously, if $m_q^2 \ll m_\chi^2$; an important new result was that this term is strongly suppressed if $m_q \geq m_\chi$, which might well be the case for the top quark. In addition we found a term $\propto (a_{\tilde{q}_i}^2 + b_{\tilde{q}_i}^2) \bar{\chi}\partial_\mu\gamma_\nu\chi F_\rho^\mu F^{\nu\rho}$. When expanded in powers of $m_{\tilde{q}}^{-2}$ it only starts at $\mathcal{O}(m_{\tilde{q}}^{-4})$; however, it is nonzero even if chirality is conserved in the (s)quark sector, i.e. if $|a_{\tilde{q}_i}| = |b_{\tilde{q}_i}|$, and therefore also receives contributions from light quarks. At first we were unable to fully exploit this result, however, for two reasons. First, we did not know the matrix element $\langle N | F_\rho^\mu F^{\nu\rho} | N \rangle$. Second, the new contribution contained terms $\propto \log \frac{m_{\tilde{q}}^2 - m_\chi^2}{m_q^2}$, i.e. are infrared divergent for light quarks.

Following a tip by Ken–Ichi Hikasa we eventually realized [18] that these problems are actually related. The logarithms can be understood, and re–summed to all orders, by using running parameters in an effective LSP–quark interaction expanded to $\mathcal{O}(m_{\tilde{q}}^{-4})$; in particular, quark operators mix with gluonic operators at one–loop level. What is more, the resulting quark and gluon operators are nothing but a subset of the so–called leading twist operators that appear in analyses of deep–inelastic lepton–nucleon scattering. Once we had realized this our work was basically done, since both the matrix of anomalous dimensions (necessary to re–sum the large logarithms in case of light quarks) and the relevant hadronic matrix elements are already well known (the latter from experiment; nonperturbative effects related to long–distance QCD are absorbed in these matrix elements, thereby solving the problem of IR divergencies). Our calculation thus extends previous studies by including all orders in $m_{\tilde{q}}^{-2}$ for terms $\propto a_{\tilde{q}_i}^2 - b_{\tilde{q}_i}^2$ (which is important for t quarks, as noted above), and by including terms $\propto (a_{\tilde{q}_i}^2 + b_{\tilde{q}_i}^2) \left(\alpha_s \log \frac{m_{\tilde{q}}^2}{m_q^2} \right)^n$ for all powers n. Comparing our results with earlier papers we also identified several sign mistakes and missing factors of 2 in the literature.[‡]

Using the basic LSP–nucleon cross section we computed LSP scattering rates in a Ge detector using the standard formalism [18]. Some results are shwon in figs. 4 and 5, where we have again used (s)particle spectra as predicted by mSUGRA. The results are quite sobering: Only in a small slice of parameter space (for positive $M \leq 100$ GeV just outside the region excluded by LEP) is the rate as large as 0.1 event/(kg day), which is the approximate limit of sensitivity that the next round of experiments aims for. On the other hand, even for quite modest sparticle masses of a few hundred GeV the rate can be as small as 10^{-3} events/(kg day), see Fig.4, or even

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[‡]For each term that had been treated in earlier papers we found at least one reference that agreed with us, but usually also at least one that did not.

10^{-4} events/(kg day) (Fig.5), due to destructive interference between the exchange of the light and heavy neutral scalar Higgs bosons. In order to detect a signal at this level one would not only have to assemble several tons of detector material, cooled to millikelvin temperatures; one would also have to suppress backgrounds by another 2 or 3 orders of magnitude compared to the goal the next round of Ge detectors is aiming for. I should add here that in these mSUGRA models squark exchange, including the terms $\propto a_{\tilde{q}_i}^2 + b_{\tilde{q}_i}^2$, is not very important, since they predict $m_\chi \leq m_{\tilde{q}}/6$ for squarks of the first two generations.[§] For a heavy nucleus like Ge73 spin–dependent interactions are usually also quite small; the total interaction is therefore dominated by Higgs exchange.

5) Summary and Conclusions

The lightest supersymmetric particle, or, more specifically, a gaugino–like neutralino, makes an excellent particle DM candidate in that its relic density more or less automatically comes out in the right ballpark even in the very restrictive class of models known as minimal supergravity models. Detection of these neutralinos might prove quite difficult, however. In order to cover a significant fraction of the allowed model parameter space one needs km^2 size neutrino detectors, or direct detection experiments that are sensitive to one event per ton of detector and day. This conclusion is somewhat more pessimistic than that of earlier studies, partly since increasing lower bounds on sparticle and Higgs boson masses exclude models with large LSP–nucleon scattering rates and partly because we insist that the LSP should have a sizable relic density to be considered a viable DM candidate. Previous studies often found large detection rates for higgsino–like or mixed LSPs, assuming a fixed local density, and ignoring the fact that such an LSP would tend to have a rather low overall relic density. Moreover, at least for the heavy top quark now favored by experiments, supergravity models predict the LSP to be gaugino–like.

One should also be aware that particle DM searches test not only particle physics but also galaxy formation models, which are necessary to estimate the density and velocity (hence the flux) of LSPs in the vicinity of the solar system. Moreover, if a signal is detected it is not clear whether it can be established as being due to superparticles (as opposed to, say, massive neutrinos). For this reason it should be clear that DM searches can never replace collider searches for supersymmetry; they might, however, allow us to "see" for the first time the stuff that most of the Universe is made of.

[§]Technically this inequality follows from the assumption of unified gaugino masses at scale M_X, as well as the fact that first generation squarks cannot be much lighter [7] than gluinos, since at one–loop level the gluino mass gives a positive contribution to the squark mass.

Acknowledgements

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Figure Captions

Fig.1 Example of contours of constant $\Omega_\chi h^2$ in the $(M, \tan\beta)$ plane for $m_t = 140$ GeV and $m = 250$ GeV (1: solid lines; 0.25: long dashed). The region outside the outer, dotted lines is excluded by various experimental and theoretical constraints other than the DM relic density.

Fig.2 Reduction of the signal of energetic muon neutrinos from LSP annihilation in the sun due to QCD effects. In the light (dark) shaded region the signal is reduced by at least 10 (50) %. The reduction is very small once $m_\chi > m_t = 150$ GeV.

Fig.3 The steps from LSP–quark scattering via Higgs and $\tilde{q}$ exchange (left) to an effective LSP–quark interaction (centre) and its connection to the gluonic matrix element $\propto F_{\mu\nu}F^{\mu\nu}$ (right). Note that in the last step the squark propagator appears inside the loop.

Fig.4 Contours of constant counting rate in a 76 Ge detector in mSUGRA. The central region between the dotted lines is excluded experimentally, mostly due to LEP chargino and Higgs searches. The solid line is the contour of 0.1 events/(kg day); in the narrow wedges between the dot–dashed lines $\Omega_\chi h^2 < 0.05$.

Fig.5 The $\tan\beta$ dependence of the LSP counting rate in a 76 Ge detector in mSUGRA. The minimum at $\tan\beta=5$ is due to destructive interference between the two Higgs exchange diagrams.

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